Heavy Baryons and electromagnetic decays

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In this talk I review the theory of electromagnetic decays of the ground state baryon multiplets with one heavy quark calculated using Heavy Hadron Chiral Perturbation Theory [1]. The M1 and E2 amplitudes for $S^* \to S\gamma$, $S^* \to T\gamma$ and $S \to T\gamma$ are separately analyzed. All M1 transitions are calculated up to $\mathcal{O}(1/\Lambda_\chi^2)$. The E2 amplitudes contribute at the same order for $S^* \to S\gamma$, while for $S^* \to T\gamma$ they first appear at $\mathcal{O}(1/(m_Q\Lambda_\chi^2))$ and for $S \to T\gamma$ are completely negligible. Once the loop contributions is considered, relations among different decay amplitudes are derived. In ref. [1] it is shown that the coupling of the photon to light mesons is responsible of a sizable enhancement of these decay widths. Furthermore, one can obtain an absolute prediction for $\Gamma(\Xi_c^{0'(*)} \to \Xi_c^0\gamma)$ and $\Gamma(\Xi_b^{-'(*)} \to \Xi_b^-\gamma)$.

1. Introduction

In Heavy Hadron Chiral Perturbation Theory (HHCPT) one constructs an effective Lagrangian whose basic fields are heavy hadrons and light mesons [2]-[5]. In ref. [6], the formalism is extended to include also electromagnetism. In this talk I describe how, using this formalism, one can calculate the electromagnetic decay width of some baryons containing a c or a b quark. The details of this computation are reported in ref.[1] and here I limit myself to trace its guidelines. In order to classify these baryons one observe that the light degrees of freedom in the ground state of a baryon with one heavy quark can be either in a $s_l = 0$ or in a $s_l = 1$ configuration. The first one corresponds to $J^P = \frac{1}{2}^+$ baryons, which are annihilated by $T_i(v)$ fields which transform as a $\bar{\mathbf{3}}$ under the chiral $SU(3)_{L+R}$ and as a doublet under the HQET $SU(2)_v$. In the second case, $s_l = 1$, the spin of the heavy quark and the light degrees of freedom combine together to form $J^P = 3/2^+$ and $J = 1/2^+$ baryons which are degenerate in mass in the $m_Q \to \infty$ limit. The spin-$\frac{3}{2}$ ones are annihilated by the Rarita-Schwinger field $S_\mu^{*ij}(v)$ while the spin-$\frac{1}{2}$ baryons are destroyed by the Dirac field $S^{ij}(v)$. They transform as a $\mathbf{6}$ under $SU(3)_{L+R}$ and as a doublet under $SU(2)_v$ and are symmetric in the i, j indices. I consider the decays $S^* \to S\gamma$ and $S^{(*)} \to T\gamma$. For most of these decays the available phase space is small, so that the emission of a pion is suppressed or even forbidden and the electromagnetic process becomes relevant. Moreover these kinds of decays are getting measured [7]. In the case of $S^* \to S\gamma$ all contributions up to order $\mathcal{O}(1/\Lambda_\chi^2)$ are calculated for M1 and E2 transitions. All divergences and scale dependence can be absorbed in the redefinition of one $\mathcal{O}(1/\Lambda_\chi)$ coupling for each type of process (M1, E2). Eliminating the unknown constants it is possible to find relations among the amplitudes which are valid up to the considered order. An analogous calculation can be performed for $S^* \to T\gamma$. In this case, the E2 contribution has to be computed up to order $\mathcal{O}(1/m_Q\Lambda_\chi^2)$, implying the intervention of two new constants. Finally for $S \to T\gamma$ the M1 amplitude is calculated up to order $\mathcal{O}(1/\Lambda_\chi^2)$, while the E2 contribution is found to be extremely suppressed. In the case $S^{(*)} \to T\gamma$ it exists a process which do not receive any contribution from local terms in the La-

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grangian and therefore its width is described by a finite chiral loop calculation: $\Gamma(\Xi_c^{0'(*)} \to \Xi_c^0\gamma)$ (and analogously $\Gamma(\Xi_b^{-'(*)} \to \Xi_b^-\gamma)$). In the following I comment these results and I refer to ref. [1] for the formalism and for a more complete comparison with other results existing in literature. A similar formalism can be applied to the study of the magnetic moments of the same baryons [8].

2. Results for $S^* \to S\gamma$ decays

The decay amplitudes are decomposed by

$$\mathcal{A}\,(B^* \to B\ \gamma)\ =\ A_{M1}\,\mathcal{O}_{M1} + A_{E2}\,\mathcal{O}_{E2}\,, \qquad (1)$$

where the corresponding M1 and E2 operators are defined by

$$\begin{aligned}\mathcal{O}_{M1} &= e\,\bar{B}\gamma_\mu\gamma_5 B^*_\nu\ F^{\mu\nu},\\ \mathcal{O}_{E2} &= i\,e\,\bar{B}\gamma_\mu\gamma_5 B^*_\nu\ v_\alpha\,(\partial^\mu F^{\alpha\nu} + \partial^\nu F^{\alpha\mu})\,, \quad (2)\end{aligned}$$

The leading contributions to M1 transitions come from the light– and heavy–quark magnetic interactions which are of $\mathcal{O}(1/\Lambda_\chi)$ and $\mathcal{O}(1/m_Q)$, respectively. We have computed the next-to-leading chiral corrections of $\mathcal{O}(1/\Lambda_\chi^2)$, which originate from the loop diagrams shown in fig. 1.

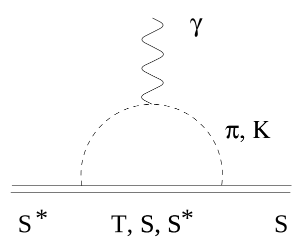

Figure 1. Meson loops contributing to $S^* \to S\gamma$.

The resulting M1 amplitudes can be written as:

$$\begin{aligned}A_{M1}(B^*) \;=\; \frac{1}{\sqrt{3}}\Bigg(&-\frac{Q_Q}{m_Q} - \frac{2c_s}{3\Lambda_\chi}\,a_\chi(B^*)\\ &+g_2^2\,\frac{\Delta_{ST}}{4(4\pi f_\pi)^2}\,a_{g_2}(B^*)\\ &+g_3^2\,\frac{m_K}{4\pi f_\pi^2}a_{g_3}\,(B^*)\Bigg)\ . \qquad (3)\end{aligned}$$

c quark	a_χ	a_{g_2}
$\Sigma_c^{++*} \to \Sigma_c^{++}\gamma$	2	$I_\pi + I_K$
$\Sigma_c^{+*} \to \Sigma_c^{+}\gamma$	1/2	$I_K/2$
$\Sigma_c^{0*} \to \Sigma_c^{0}\gamma$	-1	$-I_\pi$
$\Xi_c^{0'*} \to \Xi_c^{0'}\gamma$	-1	$-(I_\pi + I_K)/2$
$\Xi_c^{+'*} \to \Xi_c^{+'}\gamma$	1/2	$I_\pi/2$
$\Omega_c^{0*} \to \Omega_c^{0}\gamma$	-1	$-I_K$

Table 1
Contributions to M1 amplitudes for $S^* \to S\gamma$. The values of a_{g_3} can be deduced from the ones of a_{g_2} with the substitution $I_i \to m_i/m_K$ ($i = \pi, K$).

In Table 1 we show the values of the coefficients $a_i(B^*)$ for the decays of baryons containing one charm or bottom quark. In the table,

$$\begin{aligned}I_i \;\equiv\; I(\Delta_{ST}, m_i) = 2\left(-2 + \log\frac{m_i^2}{\mu^2}\right)&\\ +2\,\frac{\sqrt{\Delta_{ST}^2 - m_i^2}}{\Delta_{ST}}&\\ \times\log\left(\frac{\Delta_{ST} + \sqrt{\Delta_{ST}^2 - m_i^2}}{\Delta_{ST} - \sqrt{\Delta_{ST}^2 - m_i^2}}\right)&\ . \qquad (4)\end{aligned}$$

where Δ_{ST} is the mass difference between S and T–baryons. Due to flavor symmetry, all contributions are equal for charm and bottom baryons, with the only exception of the term proportional to the heavy quark electric charge ($Q_c = +2/3$, $Q_b = -1/3$). The main things to be observed are the following:

- the corrections proportional to g_2^2 are obtained performing a one–loop integral (fig. 1 with an S baryon running in the loop) that has to be renormalized. It can be demonstrated [1] that the scale μ dependence of the loop integrals is exactly canceled by the corresponding dependence of the coefficient $c_S(\mu)$;
- the contribution proportional to g_3^2 involves a loop integral with a baryon of the T mul-

tiplet running in the loop. Since the Lagrangian does not have any mass term for T baryons, the result of the integral is convergent and proportional to the mass of the light mesons.

Looking at table 1 one sees that relations among the decay amplitudes in which all unknown constants are eliminated can be easily found. A complete list of them is reported in ref. [1].

The M1 and E2 amplitudes have identical $SU(3)$ structure. The only difference is that there are no $1/m_Q$ terms contributing to E2. Therefore, one can construct for the E2 amplitudes exactly the same relations as in the M1 case.

The E2 amplitudes come at higher chiral order with respect to the M1 ones. Therefore, the E2 contribution to the total width is suppressed by a factor $(E_\gamma/\Lambda_\chi)^2 \sim 5\%$. In principle, it should be possible to determine experimentally the ratio A_{E2}/A_{M1} by studying the angular distribution of photons from the decay of polarized baryons [9–11]. The Fermilab E-791 experiment has reported [12] a significant polarization effect on the production of Λ_c baryons, which perhaps could be useful in future measurements of these electromagnetic decays. In ref. [1] it has been observed also that the loop contribution can strongly enhance the decay widths. In other words the coupling of the photon to light meson can give the main contribution to the decay widths.

3. Results for $S^* \to T\gamma$ decays

The M1 and E2 operators for these decays are defined as in Eq. (2). Similarly to what we have done in the previous paragraph, we write the M1 amplitude for $S^* \to T\gamma$ decays as

$$\begin{aligned} A_{M1}(B^*) &= -\sqrt{2}\,\frac{c_{ST}}{\Lambda_\chi}\,a_\chi(B^*) \\ &\quad + g_2\,g_3\,\frac{\Delta_{ST}}{2\sqrt{2}(4\pi f_\pi)^2}\,a_g(B^*) \ . \end{aligned} \tag{5}$$

The value of the parameters entering this equation can be found in ref. [1]. The final result do not depend on the heavy quark mass or charge. All constants can be eliminated in the relations

$$\begin{aligned} A_{M1}(\Sigma_c^{+*}) - A_{M1}(\Xi_c^{+'*}) &= -3\,A_{M1}(\Xi_c^{0'*}) \ , \\ A_{M1}(\Sigma_b^{0*}) - A_{M1}(\Xi_b^{0'*}) &= -3\,A_{M1}(\Xi_b^{-'*}) \end{aligned} \tag{6}$$

It is interesting to notice that $A_{M1}(\Xi_c^{0'*})$ does not depend on c_{ST}. Since at $\mathcal{O}(1/\Lambda_\chi^2)$ this decay does not get any contribution from local terms, its M1 amplitude results from a *finite* chiral loop calculation (it cannot be divergent because there is no possible counter-term to renormalize it), so that we have an absolute prediction for its value in terms of g_2 and g_3. Using the experimental value of g_3 [13,15] and the corresponding value of g_2 [16] derivable from the quark model, one finds (see also ref. [17])

$$\begin{aligned} \Gamma_{M1}(\Xi_c^{0'*}) &= 5.1 \pm 2.7 \text{ keV} \\ \Gamma_{M1}(\Xi_b^{-'*}) &= 4.2 \pm 2.4 \text{ keV} \end{aligned} \tag{7}$$

where the dominant error come from the uncertainty on $g_{2,3}$.

The E2 amplitude in $S^* \to T\gamma$ is suppressed by an extra power of $1/m_Q$. The first non-zero contributions comes at $\mathcal{O}(1/m_Q\Lambda_\chi^2)$. It is important to note that at this order it appears an operator, which break spin symmetry,

$$\begin{aligned} \mathcal{L}' &= i\,\frac{g'}{m_Q}\,\Big[\epsilon_{ijk}\,\bar{T}^i\sigma^{\mu\nu}(\xi_\mu)^j_l\,S^{kl}_\nu \\ &\quad + \epsilon^{ijk}\,\bar{S}^\mu_{kl}\sigma_{\mu\nu}(\xi^\nu)^l_j T_i\Big] \ , \end{aligned} \tag{8}$$

which gives rise to divergent loop diagrams. Moreover finite contributions of the same order come from

$$-\,i\,\frac{c_T^{E2}}{m_Q\Lambda_\chi^2}\,\epsilon_{ijk}\,\bar{T}^i\sigma_{\mu\nu}Q^j_l S^{kl}_\alpha\,\partial^\alpha\tilde{F}^{\mu\nu} \ . \tag{9}$$

Both the contributions coming from eq. 8–9 where not considered before in literature.

By eliminating the unknown coupling constants, one can deduce the relation

$$A_{E2}(\Sigma_c^{+*}) - A_{E2}(\Xi_c^{+'*}) = -3\,A_{E2}(\Xi_c^{0'*}) \ . \tag{10}$$

The same relation holds for the corresponding b baryons, since

$$A_{E2}(B_b^*) = \frac{m_c}{m_b}\,A_{E2}(B_c^*) \ . \tag{11}$$

The decays $\Xi_c^{0*} \to \Xi_c^0\gamma$ and $\Xi_b^{-*} \to \Xi_b^-\gamma$ do not get any contribution from the local term proportional to c_T^{E2}; their $O(1/m_Q\Lambda_\chi^2)$ E2 amplitude

is also given by a finite loop calculation. Unfortunately, since the coupling g' is not known, there is no absolute prediction in this case. An experimental measurement of these E2 amplitudes would provide a direct estimate of g'.

4. Results for $S \to T\gamma$

The calculation of the M1 amplitude for $S \to T\gamma$ decays is analogous to that of the previous section. Now the M1 operator is defined as

$$\mathcal{O}_{M1} = ie\,\bar{B}_T \sigma_{\mu\nu} B_S\, F^{\mu\nu} \tag{12}$$

and the corresponding amplitude can be written in the form

$$A_{M1}(B) = \frac{1}{\sqrt{6}} \frac{c_{ST}}{\Lambda_\chi}\, a_\chi(B) - g_2 g_3\, \frac{\Delta_{ST}\, a_g(B)}{4\sqrt{6}(4\pi f_\pi)^2} \;, \tag{13}$$

where the coefficients satisfy

$$a_\chi(B) = a_\chi(B^*), \qquad\qquad a_g(B) = a_g(B^*) \;. \tag{14}$$

Therefore, the relation (6) is also valid in this case. The widths of the decays $\Xi_c^{0'} \to \Xi_c^0\gamma$ and $\Xi_b^{-'} \to \Xi_b^-\gamma$ can be predicted through a finite loop calculation. From

$$\Gamma(S \to T\gamma) = 16\alpha_{em}\, \frac{E_\gamma^3 M_T}{M_S}\, |A_{M1}|^2 \;, \tag{15}$$

we find

$$\begin{aligned} \Gamma(\Xi_c^{0'}) &= (1.2 \pm 0.7)\ \mathrm{KeV} \;, \\ \Gamma(\Xi_b^{-'}) &= (3.1 \pm 1.8)\ \mathrm{KeV} \;. \end{aligned} \tag{16}$$

Again the dominant error in Eq. (16) is given by the uncertainty of $g_{2,3}$.

For these decays the E2 amplitude is further suppressed than in the previous cases. The lowest–order contribution appears at $\mathcal{O}(1/m_Q^3 \Lambda_\chi^2)$ and, therefore, can be neglected.

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